

Persistent Rabi oscillations probed via low-frequency noise correlation

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The qubit Rabi oscillations are known to be non-decaying (though with a fluctuating phase) if the qubit is continuously monitored in the weak-coupling regime. In this paper we propose an experiment to demonstrate these persistent Rabi oscillations via low-frequency noise correlation. The idea is to measure a qubit by two detectors, biased stroboscopically at the Rabi frequency. The low-frequency noise depends on the relative phase between the two combs of biasing pulses, with a strong increase of telegraph noise in both detectors for the in-phase or anti-phase combs. This happens because of self-synchronization between the persistent Rabi oscillations and measurement pulses. Almost perfect correlation of the noise in the two detectors for the in-phase regime and almost perfect anticorrelation for the anti-phase regime indicates a presence of synchronized persistent Rabi oscillations. The experiment can be realized with semiconductor or superconductor qubits.

The puzzle of the quantum state collapse due to measurement [1] is becoming accessible for the experimental study in solid state systems. Three experiments on non-projective collapse [2, 3] have been recently realized with superconducting qubits. These experiments (as well as the experiment proposed in the present paper) touch upon the most intriguing property of quantum measurement: the presence of a “spooky” quantum back-action [4], which changes the system to agree with the observation, and cannot be explained in a realistic way, i.e. by using the Schrödinger equation.

The quantum coherent (Rabi) oscillations in solid-state qubits are usually measured in an ensemble-averaged way [5, 6] and decay within a short timescale, even though it can be much longer than the oscillation period. However, for a continuous weak measurement of a single qubit, the Rabi oscillations are non-decaying and can in principle be monitored in real time, as follows, e.g., from the quantum Bayesian formalism [7], which is generally similar to the formalism of quantum trajectories [8]. Persistence of the Rabi oscillations in this case is due to the quantum back-action, which tends to increase the amplitude of the oscillations to 100%, thus competing against decoherence. The persistent Rabi oscillations lead to the spectral peak of the detector signal at the Rabi frequency [9, 10], which has been recently observed experimentally [3] (see also [11]). In the present paper we will discuss another way of demonstrating these oscillations.

For definiteness let us discuss a “charge” qubit made of a double quantum dot (DQD) populated by a single electron, the location of which is continuously measured by a nearby quantum point contact (QPC). Analogous setups can be realized with spin-based or superconducting qubits. The continuous qubit evolution due to the quantum “informational” back-action can in principle be verified in a direct experiment [7]; however, it would require high-bandwidth recording of the detector signal (including shot noise) and fast qubit manipulation, that is still a big challenge for a real experiment. A simpler way to study the back-action is to measure the qubit

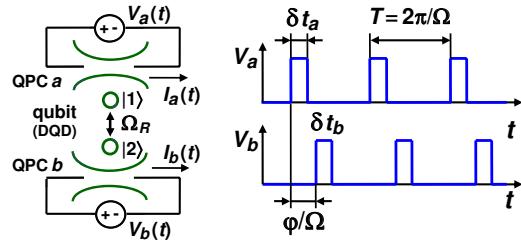


FIG. 1: Analyzed system: a double-quantum-dot qubit measured by two QPC detectors, which are biased by combs of short voltage pulses with frequency Ω coinciding with the Rabi frequency Ω_R .

by two detectors [12], so that the first (short) measurement causes a partial collapse of the qubit state, and then after a controllable qubit evolution the second detector measures the resulting state. Performing the experiment many times and selecting a certain result of the first measurement, it is possible to find the back-action evolution experimentally and compare it with the theory. The same idea with a different post-processing (selecting the result of the second measurement) has been recently used to propose an experiment on weak values [13].

The proposal of Ref. [12] still suffers from very weak signals produced by two single-shot measurements. An obvious way to increase the signal is to average it over a long comb of measurement pulses, but in this case the selection of a certain result becomes impossible. Fortunately, there is a way to overcome this dilemma by combining the ideas of two-detector measurement [12], persistent Rabi oscillations [7, 9, 10], and stroboscopic quantum non-demolition (QND) measurement [14, 15].

We propose to use the following setup (Fig. 1). A qubit is measured by two detectors, which are biased with two combs of short pulses, so that between the pulses the qubit undergoes free evolution due to the Rabi oscillations. The frequency of pulses Ω coincides with the Rabi frequency Ω_R (one pulse per period in each detector) to realize the QND regime [15]. When the two combs are

not shifted in time relative to each other, the phase of the Rabi oscillations is attracted to one of the two stable values, corresponding to the qubit being in either localized state $|1\rangle$ or $|2\rangle$ at the time of measurement. This happens because of the usual collapse in the QND frame [14, 15] and is somewhat similar to what happens with a parametrically excited swing. However, because of various imperfections (extra decoherence, etc.) there will be switching between the two stable regimes, which leads to the telegraph noise in the currents through both detectors. Even if the experimental measurement bandwidth is not wide enough to resolve the switching events (which is likely for a present-day experiment), the telegraph noise is measurable at low frequency via its spectral density, which greatly exceeds the shot noise.

The telegraph noise originates from the presence of two quasi-stable regimes of oscillations because of the QND measurement. However, for a significant phase difference φ between the two combs of the measurement pulses, the measurement is no longer QND, and the telegraph noise disappears, so that the low frequency noise reduces to a much smaller shot noise. For $\varphi \approx \pi$ the QND regime is restored again and the telegraph noise reappears. So the telegraph noise is maximum at both $\varphi = 0$ and $\varphi = \pi$; however, for $\varphi = 0$ it is almost fully correlated [16] in the two detectors (because in the stable regime both detectors see the same qubit state: either $|1\rangle$ or $|2\rangle$), while for $\varphi = \pi$ the detector noises are almost fully anticorrelated (because when one detector measures state $|1\rangle$, the other detector measures $|2\rangle$ half a period later). Experimental observation of such noise correlation/anticorrelation would show the presence of persistent Rabi oscillations.

For the quantitative analysis let us use the quantum Bayesian formalism [7] and represent the QPC currents $I_a(t)$ and $I_b(t)$ as ($n = a$ or $n = b$)

$$I_n(t) = [I_{0,n} + (\Delta I_n/2)z(t)]f_n(t) + \xi_n(t)\sqrt{|f_n(t)|}, \quad (1)$$

where $z = \text{Tr}(\sigma_z\rho)$ is the measured z -component of the qubit Bloch vector, $f_n(t)$ is the dimensionless shape of the comb of measurement pulses for n th detector (f is proportional to the QPC voltage), $I_{0,n}$ and ΔI_n are the detector average current and response for $f_n = 1$, and $\xi_n(t)$ is the shot noise with the (one-sided) spectral density S_n . We assume zero temperature. For rectangular measurement pulses we use $f_a(t) = 1$ if $|t - lT| < \delta t_a/2$ and $f_a(t) = 0$ otherwise, where δt_a is the pulse duration, $T = 2\pi/\Omega$ is the comb period, and l is an integer. Similarly, $f_b(t) = 1$ if $|t - lT - \varphi/\Omega| < \delta t_b/2$; this describes the comb of pulses of duration δt_b with the same frequency Ω but shifted by the phase φ . The qubit Hamiltonian $H_{\text{qb}} = (\Omega_R/2)\sigma_x$ describes Rabi oscillations with the frequency Ω_R about the x -axis, and we also assume pure dephasing of the qubit (not related to the measurement) with rate γ . In order to use the Markovian approximation, we assume sufficiently high QPC voltages during the pulses [7], and also assume $|\Delta I_n| \ll |I_{0,n}|$.

The spectral densities of two detector noises $S_{aa}(\omega)$ and $S_{bb}(\omega)$ as well as the cross-correlation noise $S_{ab}(\omega)$ can be found via the Fourier transform $S_{nm}(\omega) = 2\int_{-\infty}^{\infty} K_{nm}(\tau, t_0) e^{-i\omega\tau} d\tau$, where $K_{nm}(\tau, t_0) = \langle I_m(t_0 + \tau)I_n(t_0) \rangle - \langle I_m(t_0 + \tau) \rangle \langle I_n(t_0) \rangle$ is the correlation function and the averaging $K_{nm}(\tau, t_0)$ is over t_0 within one period T ; the averaging is necessary because of periodic time dependence $f_n(t)$ in Eq. (1). To find $K_{nm}(\tau, t_0)$ for $\tau > 0$ we use the method developed in [10], which essentially follows from the quantum regression theorem [17]. Expressing this correlator as $K_{nm}(\tau, t_0) = (\Delta I_n \Delta I_m / 4) f_m(t_0 + \tau) f_n(t_0) K^{zz}(\tau, t_0)$, we calculate the operator-symmetrized zz -correlator $K^{zz}(\tau, t_0)$ as

$$K^{zz}(\tau, t_0) = \sum_{i=1,2} \langle i | \rho(t_0) | i \rangle \langle i | \sigma_z | i \rangle \text{Tr}[\sigma_z \rho^{|i\rangle, t_0}(t_0 + \tau)], \quad (2)$$

where $\rho(t_0)$ is the qubit density matrix at time t_0 [$\langle i | \rho(t_0) | i \rangle = 1/2$ because of the symmetry], while $\rho^{|i\rangle, t_0}(t_0 + \tau)$ is the density matrix at time $t_0 + \tau$ for the qubit starting at time t_0 in the state $|i\rangle$, which is an eigenstate of σ_z . Here we should assume the ensemble-averaged qubit evolution, for which the measurement process is represented by pure dephasing with rates $f_n(t)\gamma_n^{\text{meas}}$, where $\gamma_n^{\text{meas}} = (\Delta I_n)^2/4S_n$. Notice that the same technique works also for two detectors measuring different observables: one of them defines the starting eigenstates, and the other one enters into the trace. While so far $\tau > 0$ was assumed, we find the correlation function at negative τ using the symmetry $K_{nm}(\tau, t_0) = K_{mn}(-\tau, t_0 + \tau)$, and also add the shot noise contribution $\delta_{nm}\delta(\tau)f_n(t_0)S_n/2$ near $\tau = 0$. Besides this quantum method, we also use below the language of a simple semiclassical analysis.

Let us first assume short measurement pulses, $\delta t_{a,b} \ll T = 2\pi/\Omega$, exactly matched frequency, $\Omega = \Omega_R$, almost no phase shift, $|\varphi| \ll 1$, and almost negligible extra dephasing, $\gamma \ll \Omega_R$. Then we have usual stroboscopic QND measurement [14, 15] insensitive to the free evolution, and therefore the qubit state eventually collapses to either $|1\rangle$ or $|2\rangle$ at the measurement moments. This obviously leads to non-decaying Rabi oscillations with 100% amplitude, which are phase-locked with the measurement combs (though with a random choice of the stable phase). The synchronization happens within the QND collapse timescale $t \sim t_{\text{col}} = T/(M_a + M_b)$, where $M_n = \gamma_n^{\text{meas}}\delta t_n = \delta t_n(\Delta I_n)^2/4S_n$. We assume $M_n \ll 1$, so that $t_{\text{col}} \gg T$, while $\gamma_{a,b}^{\text{meas}}T$ are not necessarily small. In the ideal QND case the phase of the Rabi oscillations is fixed forever after this gradual collapse; however, in a realistic case there will be switching between the two regimes (state $|1\rangle$ or $|2\rangle$ at the measurement moments) with a calculated below rate Γ_S , the same for both switching processes because of the symmetry. If we assume rare switching, $\Gamma_S t_{\text{col}} \ll 1$, then the detector current I_n averaged over a coarse graining timescale

longer than T and t_{col} , switches between the two levels, $(I_{0,n} \pm \Delta I_n/2)(\delta t_n/T)$, thus producing the telegraph noise. Therefore, the noise spectral density $S_{nn}(\omega)$ at frequencies $\omega \ll t_{\text{col}}^{-1} \ll \Omega$, is

$$S_{nn}(\omega) = \left(\frac{\delta t_n}{T}\right)^2 \frac{(\Delta I_n)^2/2\Gamma_S}{1 + (\omega/2\Gamma_S)^2} + \frac{\delta t_n}{T} S_n, \quad (3)$$

where the term $(\delta t_n/T)S_n$ is due to the shot noise. It is easy to see that at low frequency, $\omega \ll \Gamma_S$, the ratio of the telegraph and shot noise contributions $(\delta t_n/T)(\Delta I_n)^2/2\Gamma_S S_n \simeq 1/t_{\text{col}}\Gamma_S$ is always large in our case.

For the phase shift $\varphi \approx \pi$ the QND regime is still realized, and therefore Eq. (3) for each detector noise is still valid. However, since the detectors now measure the opposite qubit states, the cross-correlation noise changes sign,

$$S_{ab}(\omega) = \pm \frac{\delta t_a \delta t_b}{T^2} \frac{\Delta I_a \Delta I_b / 2\Gamma_S}{1 + (\omega/2\Gamma_S)^2}, \quad (4)$$

where “+” sign is for $\varphi \approx 0$ and “-” is for $\varphi \approx \pi$. The noise correlation factor $S_{ab}(0)/\sqrt{S_{aa}(0)S_{bb}(0)}$ is obviously close to ± 1 , describing almost full correlation/anticorrelation, when the shot noise term in Eq. (3) is much smaller than the telegraph noise.

To find the switching rate Γ_S , we calculate the “propagators” $\rho^{|i\rangle, t_0}(t_0 + \tau)$ in Eq. (2), essentially rederiving Eqs. (3) and (4) in the fully quantum way. Because of assumed weak coupling ($\gamma T \ll M_{a,b} \ll 1$) these density matrices at time $t = t_0 + \tau$ can be represented in the Bloch coordinates as $z = A(t) \cos[\Omega t - \phi(t)]$, $y = \text{Tr}(\sigma_y \rho) = A(t) \sin[\Omega t - \phi(t)]$ with slowly changing amplitude $A(t)$ and phase $\phi(t)$:

$$\dot{A} = -(A/T)[M_a \sin^2(-\phi) + M_b \sin^2(\varphi - \phi)] - \gamma A/2, \quad (5)$$

$$\dot{\phi} = (1/2T)[M_a \sin(-2\phi) + M_b \sin(2\varphi - 2\phi)], \quad (6)$$

where we assumed $\delta t_{a,b} \ll T$, so that periodic instantaneous dephasings of magnitudes M_a and M_b happen at the phases $-\phi$ and $\varphi - \phi$. Assume now $t_0 = 0$ [so that $f_a(t_0) = 1$] and choose $|i\rangle = |1\rangle$; then $A(0) = 1$ and $\phi(0) = 0$. For $|\varphi| \ll 1$ the solution of Eq. (6) is simple, and ϕ saturates at $\phi_{st} = \varphi M_b / (M_a + M_b)$ exponentially with time constant t_{col} . This value can be inserted into Eq. (5) because evolution of the amplitude A due to measurement is much slower than t_{col}^{-1} , resulting in $A(t) = \exp[-(t/T)\varphi^2 M_a M_b / (M_a + M_b) - \gamma t/2]$. It is easy to see that the evolution starting with the state $|2\rangle$ leads to ϕ shifted by π , but the same $A(t)$, and the same contribution into $K^{zz}(\tau, 0)$ in Eq. (2). Calculating now $S_{aa}(\omega)$ via $K^{zz}(\tau, 0)$, and using approximation $|\cos \phi_{st}| \approx 1$ because $|\phi_{st}| < |\varphi| \ll 1$, we obtain the formula, coinciding with Eq. (3) with $\Gamma_S = (1/2T)\varphi^2 M_a M_b / (M_a + M_b) + \gamma/4$. Calculation of $S_{bb}(\omega)$ is fully similar, while to obtain $S_{ab}(\omega)$ in the form (4) with the same switching rate Γ_S we also need approximation $|\cos(\phi - \phi_{st})| \approx 1$.

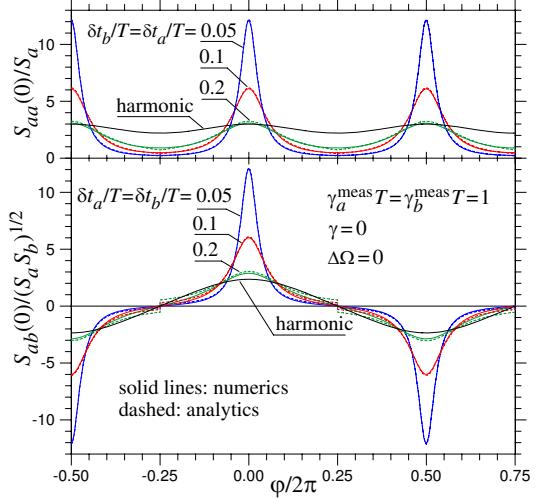


FIG. 2: (Color online) Numerical (solid lines) and analytical (dashed lines) dependence of zero-frequency detector noise $S_{aa}(0)$ and cross-noise $S_{ab}(0)$ on the phase shift φ between the bias voltage combs for several values of the pulse width $\delta t_{a,b}$, and also for the harmonic biasing. Almost complete noise anticorrelation at $\varphi = \pm\pi$ indicate persistent Rabi oscillations.

To account for small non-zero pulse widths $\delta t_{a,b}$, we can still use Eq. (6), but averaging over ϕ within the pulse widths in Eq. (5) leads to the extra factor $\exp[-(2\pi)^2(M_a \delta t_a^2 + M_b \delta t_b^2)t/12T^3]$ in $A(t)$ and corresponding increase of Γ_S . A small frequency mismatch $\Delta\Omega = \Omega - \Omega_R$ would lead to the extra term $\Delta\Omega$ in Eq. (6), so that ϕ_{st} becomes shifted by $\Delta\Omega T / (M_a + M_b)$. In the case when the shift between the measurement combs is close to the half-period, φ should be obviously replaced by $\varphi \pm \pi$. Taking into account these changes, we reproduce Eqs. (3) and (4) with the switching rate

$$\Gamma_S = \frac{\tilde{\varphi}^2 M_a M_b + (\Delta\Omega T)^2}{2T(M_a + M_b)} + \frac{M_a \delta t_a^2 + M_b \delta t_b^2}{6T^3/\pi^2} + \frac{\gamma}{4} \quad (7)$$

for $\tilde{\varphi} \ll 1$, where $\tilde{\varphi} = \min(|\varphi|, |\varphi \pm \pi|)$. Assuming comparable measurement parameters for both detectors, we see that the telegraph noise at zero frequency greatly exceeds the shot noise contribution in Eq. (3) if $\tilde{\varphi} \ll 1$, $\delta t_{a,b} \ll T$, $|\Delta\Omega|T \ll M_{a,b}$, and $\gamma T \ll M_{a,b}$. This is the condition for the validity of our analytical results.

Figure 2 shows zero-frequency spectral densities $S_{aa}(0)$ and $S_{ab}(0)$ as functions of the phase shift φ for several values of the pulse width $\delta t_{a,b}$, assuming negligible $\Delta\Omega$ and γ . Solid lines are the numerical results calculated via Eq. (2), while the dashed lines show analytical results using Eqs. (3), (4), and (7). Overall the analytics is very close to the numerical results (almost coinciding), except for $S_{ab}(0)$ near $\varphi = \pm\pi/2$, where the analytics is discontinuous because of the sign change in (4). We normalize the noise $S_{aa}(0)$ by the shot noise S_a of constantly biased detector, so that the shot noise contribution in this normalization is $\delta t_a/T$. Similarly, the cross-noise

$S_{ab}(0)$ is normalized by $\sqrt{S_a S_b}$. The numerical results in Fig. 2 confirm almost full correlation of the detector noises at $\varphi \approx 0$ and almost full anticorrelation at $\varphi \approx \pm\pi$ [16]. The peaks become higher and narrower for shorter pulse durations $\delta t_{a,b}$. The results in the used normalization are practically insensitive to the qubit-detector coupling $\gamma_{a,b}^{\text{meas}}$ (assuming $M_{a,b} \ll 1$). Non-zero detuning $\Delta\Omega$ and/or extra dephasing γ make the peaks in Fig. 2 lower, while not affecting their width; this lowering is less significant for stronger coupling $\gamma_{a,b}^{\text{ms}}$. We have also checked numerically that the frequency dependence of the noises at $\omega \ll \Omega$ is close to the analytical results (3) and (4); extra peaks as well as significant imaginary component of $S_{ab}(\omega)$ appear at $\omega \approx \Omega$ and overtones of Ω .

Obviously, the analysis and results change only trivially if Ω_R/Ω is an integer or close to an integer. In a real experiment with QPC detectors the best measurement mode is to apply two bias voltage pulses with opposite polarity per Rabi period for each detector. In this case the average bias voltage is zero that helps to keep zero bias between the pulses. The average current in each detector is then also zero, simplifying the noise measurement. For such mode $\delta t_{a,b}$ in Eqs. (3) and (4) should be replaced by $2\delta t_{a,b}$, while in Eq. (7) the measurement strengths $M_{a,b}$ should be doubled (no change for $\delta t_{a,b}$).

Now let us discuss why experimental observation of the noise dependence of Fig. 2 would indicate persistent Rabi oscillations. Correlation of the noises for $\varphi \approx 0$ could be alternatively explained by the qubit localization in either state $|1\rangle$ or $|2\rangle$. However, the anticorrelation for $\varphi \approx \pm\pi$ is possible only if the qubit oscillates persistently. Moreover, these oscillations should be synchronized with the measurement combs, because for persistent Rabi oscillations with a random phase one would expect dependence $S_{ab}(0) \propto \cos \varphi$, which is very different from the peaked dependence in Fig. 2. One may also worry that the noise dependence of Fig. 2 could be alternatively explained by the driven Rabi oscillations (between the energy eigenstates) caused by presence of a voltage with resonant frequency. However, both energy eigenstates produce no signal in the detectors; therefore the driven Rabi oscillations could only reduce the discussed noise correlation and cannot be used for an alternative explanation (notice also that both stable phases of the persistent oscillations are insensitive to the microwave drive $\propto \cos \Omega_R t$). Unfortunately, measurement of only zero-frequency noise is insufficient to demonstrate $\sim 100\%$ amplitude of the persistent Rabi oscillations (observed in [3]). However, if the switching rate Γ_S can be measured either in frequency or time domain, then by comparing $S_{ab}(0)$ with Eq. (4) one can check the amplitude of synchronized oscillations.

If the stroboscopic biasing is replaced by harmonic biasing [18]: $f_a = \cos(\Omega t)$, $f_b = \cos(\Omega t - \varphi)$, then $S_{ab}(0)$ still depends on the phase shift φ (see Fig. 2); however, there are no more peaks and the noise magnitude is relatively small. The numerical results at weak coupling can

be fitted as $S_{ab}(0) = 1.18\Delta I_a \Delta I_b \cos \varphi / (\gamma_a^{\text{meas}} + \gamma_b^{\text{meas}})$ (actually, the φ -dependence is slightly more peaked than $\cos \varphi$).

In our analysis we have neglected the noise from amplifiers, which can be simply added and is not expected to depend on φ . The main effect of the neglected thermal noise in the QPCs is a small contribution to the dephasing γ between the pulses. A weak energy relaxation in the qubit can also be easily taken into account.

For numerical estimates let us assume QPCs with $I_{a,b} \simeq 100$ nA, $\Delta I_{a,b}/I_{a,b} \simeq 0.1$, symmetric biasing with $\delta t_{a,b}/T \simeq 0.1$, and Rabi frequency $\Omega_R/2\pi \simeq 2$ GHz. Then the collapse (“attraction”) time $t_{\text{col}} \simeq 2$ ns is few Rabi periods, while the switching rate is $\Gamma_S \simeq \tilde{\varphi}^2/15$ ns + $1/120$ ns + $(\Delta\Omega/\Omega)^2/6$ ps + $\gamma/4$. Therefore we need the dephasing time $T_2 = 1/\gamma$ to be longer than only few ns to have significant correlated telegraph noise, and its ratio to the shot noise contribution for $\tilde{\varphi} = \Delta\Omega = 0$ is crudely $\min(60, T_2/0.5\text{ns})$ (5 times smaller for the normalization of Fig. 2). These figures show that the experiment is doable using the present-day semiconductor technology [6, 19]. The experiment can also be realized with the superconducting qubit setup of Ref. [3]. In comparison with the experiment [3] it would also demonstrate partial synchronization of the persistent Rabi oscillations without use of a much more complicated quantum feedback.

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